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Scale Invariant Volkov–Akulov Supergravity

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Abstract

A scale invariant Goldstino theory coupled to supergravity is obtained as a standard supergravity dual of a rigidly scale–invariant higher–curvature supergravity with a nilpotent chiral scalar curvature. The bosonic part of this theory describes a massless scalaron and a massive axion in a de Sitter Universe.

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1 Introduction

Motivated by single–field inflationary scenarios [1], several sgoldstinoless [2–5] supergravity extensions of inflationary models were recently considered [6–11] (for a recent review see [12]). Interestingly enough, in [7, 11] many of these models were linked to pure higher–derivative supergravity with a nilpotency constraint on the scalar curvature chiral superfield $\mathcal{R}$. These include the Volkov–Akulov–Starobinsky model [7] and the pure Volkov–Akulov theory coupled to supergravity [7]. Recently, the full component form of the latter theory was presented in [13, 14]

Along these lines, various authors considered R^2 theories of gravity [15] and their supergravity embeddings [15, 16], which possess a rigid scale invariance and naturally accommodate a de Sitter Universe. It is the aim of this note to give the sgoldstinoless version of these theories, which naturally combines an enhanced rigid scale invariance and a de Sitter geometry. This theory also emerges as a limiting case of the inflationary scenario.

2 Scale–Invariant Nilpotent Supergravity

The superspace action density of the scale–invariant theory that we consider [3],

$$\mathcal{A} \ = \ \frac{\mathcal{R}\,\overline{\mathcal{R}}}{g^2}\bigg|_D \ + \ \sigma\,\mathcal{R}^2\,S_0\bigg|_F \ , \tag{2.1}$$

where g is a dimensionless parameter, is invariant under the rigid scale transformations

$$\mathcal{R} \ \to \ \mathcal{R} \ , \ S_0 \ \to \ e^{-\lambda}\,S_0 \ , \ \sigma \ \to \ e^{\lambda}\,\sigma \ . \tag{2.2}$$

This theory is equivalent to the theory considered in [16], supplemented with the nilpotency constraint

$$\mathcal{R}^2 \ = \ 0 \ , \tag{2.3}$$

which is enforced by the chiral Lagrange multiplier σ present in the second term of eq. (2.1).

Using manipulations similar to those originally introduced in [17], we can now turn this model into a scale–invariant version of the Volkov–Akulov model coupled to standard supergravity. To this end, we first use the superspace identity

$$\sigma\,\mathcal{R}^2\,S_0 \ + \ h.c.\bigg|_F \ = \ \left(\sigma\,\frac{\mathcal{R}}{S_0} \ + \ \overline{\sigma}\,\frac{\overline{\mathcal{R}}}{\overline{S}_0}\right) S_0\,\overline{S}_0\bigg|_D \ + \ \text{tot. deriv.} \ , \tag{2.4}$$

[3]We use throughout the conventions of [7].

and then introduce two Lagrange chiral superfield multipliers T and S according to

$$\mathcal{A} \;=\; \left(\sigma\, S \;+\; \overline{\sigma}\,\overline{S} \;+\; \frac{S\,\overline{S}}{g^2}\right) S_0\,\overline{S}_0\Big|_D \;-\; T\left(\frac{\mathcal{R}}{S_0} \;-\; S\right) S_0^3 \;+\; \text{h.c.}\Big|_F \;. \tag{2.5}$$

The final result is the standard supergravity action density

$$\mathcal{A} \;=\; -\,\left(T + \overline{T} - \sigma\, S - \overline{\sigma}\,\overline{S} \;-\; \frac{S\,\overline{S}}{g^2}\right) S_0\,\overline{S}_0\Big|_D \;+\; T\,S\,S_0^3 \;+\; \text{h.c.}\Big|_F \;+\; \text{tot. deriv.} \tag{2.6}$$

A final shift and a redefinition according to

$$T \;\to\; T \;+\; \sigma\, S \;, \qquad X \;=\; \frac{S}{g} \tag{2.7}$$

yield the standard supergravity action density

$$\mathcal{A} \;=\; -\,(T \;+\; \overline{T} \;-\; X\,\overline{X})\, S_0\,\overline{S}_0\Big|_D \;+\; W(T,X)\, S_0^3 \;+\; \text{h.c.}\Big|_F \;, \tag{2.8}$$

where

$$W(T,X,\sigma) \;=\; g\,T\,X \;+\; g^2\,\sigma\, X^2 \;. \tag{2.9}$$

This is tantamount to the scale–invariant superpotential

$$W(T,X) \;=\; g\,T\,X \;, \tag{2.10}$$

where X is subject to the nilpotency constraint

$$X^{\,2} \;=\; 0 \;, \tag{2.11}$$

so that X describes the sgoldstinoless Volkov–Akulov multiplet [2–5]. The corresponding bosonic Lagrangian,

$$\mathcal{L} \;=\; \frac{R}{2} \;-\; \frac{3}{(T \,+\, \overline{T})^2}\;|\partial\, T|^2 \;-\; g^{\,2}\,\frac{|T|^2}{3(T \,+\, \overline{T})^2} \;, \tag{2.12}$$

is a special case of the result displayed in [7], so that it describes an $SU(1,1)/U(1)$ Kählerian model of curvature -2/3 with a scale–invariant positive potential. As a result, in terms of the canonical variable

$$T \;=\; e^{\phi\sqrt{\frac{2}{3}}} \;+\; i\,a\,\sqrt{\frac{2}{3}} \;, \tag{2.13}$$

one finds

$$\mathcal{L} \;=\; \frac{R}{2} \;-\; \frac{1}{2}\,(\partial\phi)^2 \;-\; \frac{1}{2}\; e^{\,-\,2\,\phi\,\sqrt{\frac{2}{3}}}\;(\partial a)^2 \;-\; \frac{g^{\,2}}{12} \;-\; \frac{g^{\,2}}{18}\; e^{-2\,\phi\sqrt{\frac{2}{3}}}\; a^2 \;. \tag{2.14}$$

Note that in the Einstein frame the metric is inert under the scale transformation corresponding to eq. (2.2), while

$$\phi \;\to\; \phi \;+\; \gamma \;, \quad a \;\to\; e^{\gamma\,\sqrt{\frac{2}{3}}}\; a \;. \tag{2.15}$$

3 de Sitter Vacuum Geometry

Since a is stabilized at zero, this model results in a de Sitter vacuum geometry, with a corresponding scale–invariant realization of supersymmetry breaking induced by the non–linear sgoldstinoless multiplet. The supersymmetry breaking scale M_s^2 is

$$M_s^2 \;=\; \frac{g}{2\sqrt{3}}\, M_{Planck}^2 \;, \tag{3.1}$$

up to a conventional numerical factor. Eq. (2.8) describes the minimal supergravity model that embodies a scale–invariant goldstino interaction and leads unavoidably to a de Sitter geometry. This model involves a single dimensionless parameter g, which determines its *positive* vacuum energy according to

$$V \;=\; \frac{g^2}{12}\, M_{Planck}^4 \;. \tag{3.2}$$

In contrast, the Volkov–Akulov model coupled to supergravity, depends on the two parameters f and W_0, and consequently leads to a vacuum energy [18] [19] [20] [7] [13]

$$V \;=\; \frac{1}{3}\,|f|^2 \;-\; 3\,|W_0|^2 \tag{3.3}$$

of arbitrary sign.

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